

Higgs-dependent Yukawa couplings

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Abstract

We consider the possibility that the Yukawa couplings depend on the Higgs field, with the motivation of generating the fermion mass hierarchy through appropriate powers of the Higgs vacuum expectation value. This leads to drastic modifications of the Higgs branching ratios, new Higgs contributions to various flavor-violating processes, and observable rates for the top quark decay $t \rightarrow hc$. The underlying flavor dynamics must necessarily appear at the TeV scale and is within the reach of the LHC.

1 Introduction

The theoretical understanding of fermion masses and mixing angles is one of the most important unresolved problems in particle physics. In this paper we will try to address this problem by introducing an effective theory, valid below a new mass scale M , in which the Yukawa couplings depend upon the Higgs field H . The theory does not contain any small coupling constants and the hierarchical pattern of the fermion masses is understood in terms of powers of $\langle H \rangle / M$. The new mass scale M turns out to be necessarily around the TeV, suggesting a possible link between flavor dynamics and the physics associated with electroweak symmetry breaking. A similar proposal was made in ref. [1].

At first sight it may seem that introducing flavor dynamics at the TeV scale will have disastrous consequences for flavor-changing neutral current (FCNC) processes. Indeed, in the effective theory, new contributions to FCNC are already induced at tree level by Higgs-boson exchange. Nevertheless, as we will show in sect. 3, the effect is sufficiently small, since the flavor-violating Higgs couplings are related to the pattern of fermion masses. An explicit example of how Higgs-dependent Yukawa couplings can be generated at the scale M is presented in sect. 5.

Since the scale of flavor dynamics M is linked to the electroweak scale, our approach leads to many experimentally-observable distinctive features. New flavor-violating effects are expected just beyond the present experimental bounds. The new states at the scale M are well within the reach of the LHC. Moreover, as we describe in sect. 4, the Higgs branching ratios are drastically modified, because the Higgs couplings to fermions are larger than those in the Standard Model (SM) by a factor of a few. Finally, the Higgs boson can decay in flavor-violating channels or, depending on its mass, be produced in the top-quark decay $t \rightarrow hc$.

2 The effective theory of Higgs-dependent Yukawa couplings

Our basic assumption is that, in an effective theory valid below the new-physics scale M , quark and lepton Yukawa couplings $Y^{u,d,\ell}$ are functions of the Higgs field H . Therefore, the Yukawa couplings can be expanded as

$$Y_{ij}(H) = \sum_{n=0}^{\infty} c_{ij}^{(n)} \left(\frac{H^\dagger H}{M^2} \right)^n, \quad (1)$$

where i, j are generation indices. In what follows we focus upon the particularly interesting possibility that the hierarchical pattern of fermion masses is explained by appropriate powers of $\langle H \rangle / M$. Therefore, we assume that the coefficients $c_{ij}^{(n)}$ vanish up to a (generation-dependent) order n_{ij} . Keeping only the leading order terms in eq. (1), we can express the Yukawa effective Lagrangian as

$$-\mathcal{L}_Y = Y_{ij}^u(H) \bar{q}_{Li} u_{Rj} H^c + Y_{ij}^d(H) \bar{q}_{Li} d_{Rj} H + \text{h.c.} , \quad Y_{ij}^{u,d}(H) = c_{ij}^{u,d} \left(\frac{H^\dagger H}{M^2} \right)^{n_{ij}^{u,d}} , \quad (2)$$

where $H^c = i\sigma_2 H^*$. Here we restrict our considerations to quarks and comment on the case of leptons at the end of sect. 3. The coefficients $c_{ij}^{u,d}$ are numbers of order unity, while $n_{ij}^{u,d}$ are integers.

When the Higgs field develops a vacuum expectation value (VEV) $\langle H \rangle \equiv v \ll M$ (where $v = 174$ GeV), the fermions attain hierarchically small masses, depending on $n_{ij}^{u,d}$. On the other hand, the coupling of the physical Higgs boson $h = \sqrt{2}\text{Re}(H^0 - v)$ to the fermions with flavors i, j increases by a factor $2n_{ij} + 1$ compared to that of the SM,

$$y_{ij}^{u,d} = \left(2n_{ij}^{u,d} + 1 \right) (y_{ij}^{u,d})_{\text{SM}} , \quad (3)$$

where $(y_{ij}^{u,d})_{\text{SM}} = m_{ij}^{u,d} / (\sqrt{2}v)$. For the top quark $n_{33}^u = 0$, while suppression of the bottom quark mass requires $n_{33}^d = 1$. This fixes the expansion parameter

$$\epsilon \equiv \frac{v^2}{M^2} \simeq \frac{m_b}{m_t} \simeq \frac{1}{60} , \quad (4)$$

and therefore the new-physics scale M must be about 1–2 TeV. It is intriguing that flavor physics points independently towards the same scale that is favored by hierarchy-problem considerations. Note that choosing $n_{33}^d > 1$ would lead to an excessively low value of M .

The structure of Yukawa couplings in eq. (2) is reminiscent of the Froggatt-Nielsen approach [2]. However, there are three important differences. First of all, in the Froggatt-Nielsen case, the new-physics scale is arbitrary and often associated with some super-heavy mass. In our case, M is necessarily around the TeV and this leads to a rich variety of phenomenological predictions that we will discuss in the following.

The second difference is that our expansion parameter ϵ in eq. (4) is smaller than the typical expansion parameter of the Froggatt-Nielsen approach, which is usually taken to be equal to the Cabibbo angle $\lambda = 0.22$. This means that we will use the ϵ expansion to reproduce only the broad features of the fermion mass matrices and allow the coefficients $c_{ij}^{u,d}$ to take values from about 1/5 to 5, in order to fit precisely the mass and mixing parameters.

In practice, we will choose appropriate values of $n_{ij}^{u,d}$ to obtain the hierarchy

$$\frac{m_t}{v} \sim V_{us} \sim \mathcal{O}(\epsilon^0), \quad \frac{m_{b,c}}{v} \sim V_{cb} \sim V_{ub} \sim \mathcal{O}(\epsilon^1), \quad \frac{m_s}{v} \sim \mathcal{O}(\epsilon^2), \quad \frac{m_{u,d}}{v} \sim \mathcal{O}(\epsilon^3). \quad (5)$$

The third important remark is that $H^\dagger H$ cannot carry any quantum number and thus cannot play the role of the Froggatt-Nielsen field. However, this becomes possible in the supersymmetric version of our approach, where one replaces $H^\dagger H$ with the gauge-invariant combination $H_u H_d$ of the two Higgs doublet superfields. The Yukawa couplings in the superpotential are holomorphic functions of this combination and can be expressed as

$$Y_{ij}^{u,d} = c_{ij}^{u,d} \left(\frac{H_u H_d}{M^2} \right)^{n_{ij}^{u,d}}. \quad (6)$$

Since the field combination $H_u H_d$ can carry a U(1) charge, it provides a direct analogy to the Froggatt-Nielsen field. The form of Y_{ij} is dictated by U(1) charge conservation and with an appropriate charge assignment one can reproduce the fermion masses and mixings. In this case, the Yukawa couplings have a “factorizable” form and

$$n_{ij}^{u,d} = a_i + b_j^{u,d}, \quad (7)$$

where $a_i, b_j^{u,d}$ are related to the U(1) charges of q_{Li} , u_{Rj} and d_{Rj} , respectively. We also note that here the expansion parameter is $\epsilon = (v^2 \sin 2\beta)/(2M^2)$, which is smaller than that in the non-supersymmetric case for the same value of M , especially at large $\tan \beta$. However, large $\tan \beta$ also reduces the ratio between the top and bottom Yukawa couplings.

The absence of any quantum number associated with $H^\dagger H$ makes the form of the Lagrangian in eq. (2) potentially unstable under quantum corrections. Indeed, the power n_{ij} in the Yukawa interaction can be reduced by closing the $H^\dagger H$ lines through a Higgs propagator into a quadratically divergent one-loop diagram. Let us assume that all quadratic divergences associated with the Higgs are cut off by TeV-physics at the mass scale of the new states, which is equal to $g_F M$, where g_F is the typical coupling of the unknown flavor dynamics. The reduction of n_{ij} by one unit, obtained by the one-loop integration, carries a suppression factor of about $g_F^2/(16\pi^2)$. This effect is subleading to the hierarchy created by the ϵ expansion as long as

$$\frac{g_F^2}{16\pi^2} < \epsilon. \quad (8)$$

In the following, we will assume the validity of eq. (8), thus ensuring that the pattern of Yukawa couplings described by eq. (2) is rather stable under radiative corrections. Note that eq. (8) requires that flavor dynamics be rather weakly coupled ($g_F < 4\pi\sqrt{\epsilon} \simeq 1$), and

it gives an independent argument for fixing the flavor-dynamics mass scale M in the TeV range ($M < 4\pi v/g_F \simeq 2 \text{ TeV}/g_F$).

In general, the effective theory also contains Higgs-dependent fermionic kinetic terms of the form

$$\mathcal{L}_K = iZ_{ij}^q(H)\bar{q}_{Li}\not{D}q_{Lj} + \left(i\tilde{Z}_{ij}^q(H)\bar{q}_{Li}\not{D}[\alpha_{ij}(H)q_{Lj}] + \text{h.c.}\right) + \dots, \quad (9)$$

where Z^q is a Hermitian matrix depending on $H^\dagger H/M^2$, \tilde{Z}^q is a general matrix depending on H and H^\dagger , and $\alpha_{ij}(H)$ is a function of H and H^\dagger (e.g. it is either H or H^\dagger to lowest order). The ellipses stand for higher derivative terms. Similar expressions hold for u_R and d_R kinetic terms.

Consider the first term in eq. (9). When H is replaced by its VEV, the kinetic terms can be made canonical by unitary rotations and field rescalings. These transformations affect the Yukawa matrices and, if the coefficients $Z_{ij}^{q,u,d}(v)$ contain non-trivial powers of v/M , they can contribute to generation of the hierarchy in fermion masses and mixings. However, if quarks are elementary particles in the fundamental theory, then $Z^{q,u,d} = \mathbb{1} + \mathcal{O}[(H^\dagger H/M^2)^n]$ (with $n > 0$) and the field rescaling gives only subdominant effects in ϵ . The terms in $Z^{q,u,d}$ depending on the Higgs field fluctuation cannot be removed and generate additional anomalous Higgs interactions. Since we are interested in processes in which the fermions are real particles in the external legs, we can use their equations of motion to reduce these interactions to the form of Yukawa couplings. Therefore, for simplicity, in the following we will drop the effects of the kinetic terms in eq. (9) and concentrate only on the Higgs-dependent Yukawa couplings.

The second term in eq. (9) yields corrections to both the kinetic term and the Z - and W -couplings. This is because the gauge structure of $\not{D}[\alpha(H)q_L]$ is in general different from that of $\not{D}q_L$ and after electroweak symmetry breaking this produces a correction to the couplings between fermions and gauge bosons. Such a correction in general splits into a flavor universal part, which affects electroweak precision physics, and a flavor violating part which is constrained by flavor physics. These effects can be studied only on a model-by-model basis. We also note that at 1-loop additional dipole-type interactions $\bar{q}_{Li}\sigma_{\mu\nu}q_{Rj}F^{\mu\nu}$ are generated.

3 Higgs-mediated flavor violation

From eq. (3) it immediately follows that the Higgs boson mediates FCNC at tree level. The reason is that the fermion mass matrix and the matrix of Higgs-fermion couplings differ by

a flavor dependent factor $2n_{ij} + 1$, and thus they cannot be diagonalized simultaneously.

Let us choose a field basis in which the quark mass matrices $m^{u,d}$ are diagonal and real, obtained through the bi-unitary transformation

$$Y^{u,d}(v) \rightarrow V_L^{u,d\dagger} Y^{u,d}(v) V_R^{u,d} = \frac{m^{u,d}}{v}. \quad (10)$$

Here $V_{L,R}^{u,d}$ are unitary matrices and the CKM matrix is given by $V_{CKM} = V_L^{u\dagger} V_L^d$. In this basis, the interaction between a single Higgs boson and the fermionic current is given by

$$\mathcal{L}_h = -\frac{h}{\sqrt{2}} J_h, \quad J_h \equiv \frac{m_i^u}{v} \bar{u}_i u_i + 2 (G_{ij}^u \bar{u}_{Li} u_{Rj} + \text{h.c.}) + (u \leftrightarrow d), \quad (11)$$

$$G_{ij}^{u,d} \equiv \frac{m_k^{u,d}}{v} n_{lm}^{u,d} V_{L li}^{u,d*} V_{L lk}^{u,d} V_{R mk}^{u,d*} V_{R mj}^{u,d}. \quad (12)$$

The coupling $G_{ij}^{u,d}$ describes the new interaction term, not present in the SM. Its expression becomes simpler under the factorization hypothesis of eq. (7),

$$G_{ij}^{u,d} = A_{ij}^{u,d} \frac{m_j^{u,d}}{v} + \frac{m_i^{u,d}}{v} B_{ij}^{u,d}, \quad A^{u,d} \equiv V_L^{u,d\dagger} a V_L^{u,d}, \quad B^{u,d} \equiv V_R^{u,d\dagger} b^{u,d} V_R^{u,d}. \quad (13)$$

The information about flavor violation is contained in the Hermitian matrices $A^{u,d}$ and $B^{u,d}$. To simplify our discussion, we will focus on factorizable Yukawa matrices, but analogous considerations can be made for the general case.

By integrating out the Higgs boson at tree level, we obtain the four-fermion effective interaction Lagrangian

$$\mathcal{L}_{4f} = \frac{J_h^2}{4m_h^2}, \quad (14)$$

where m_h is the Higgs boson mass. The Lagrangian in eq. (2) also contains multi-Higgs interactions with quarks of the form $m_{ij} \bar{q}_i q_j (h/v)^{p+1}$, where $p \leq 2n_{ij}$ and m_{ij} is the quark mass matrix in the current eigenbasis. By integrating out the Higgs bosons at p loop order, these interactions also generate terms in the four-fermion effective Lagrangian. The ratio between the coefficients of the p -loop contribution and the tree-level contribution is of the order of

$$\left(\frac{g_F M}{4\pi v} \right)^{2(p-1)} \left(\frac{m_h}{4\pi v} \right)^2, \quad (15)$$

where we have cut off power-divergent integrals at the mass scale of the new states $g_F M$. Using eq. (8) and the requirement of a perturbative Higgs quartic coupling ($m_h < 4\pi v$), we observe that loop contributions are always subleading with respect to the tree-level effect given in eq. (14).

The flavor-violating effects can be easily extracted from eq. (14). For instance, the $\Delta S = 2$ interaction is given by

$$\mathcal{L}_{\Delta S=2} = \frac{m_s^2}{v^2 m_h^2} \left(A_{12}^d \bar{d}_L s_R + B_{12}^d \bar{d}_R s_L \right)^2. \quad (16)$$

This gives a contribution to the mass difference of the neutral kaons

$$\frac{\Delta m_K}{m_K} \simeq \frac{5 f_K^2 m_K^2}{12 v^2 m_h^2} \left(A_{12}^{d\,2} + B_{12}^{d\,2} - \frac{12}{5} A_{12}^d B_{12}^d \right). \quad (17)$$

Analogous results can be obtained for the mass differences in the B^0 and D^0 systems. Requiring that the new contributions do not exceed the experimental bounds, we obtain the following constraints

$$\sqrt{\left| A_{12}^{d\,2} + B_{12}^{d\,2} - \frac{12}{5} A_{12}^d B_{12}^d \right|} < 6 \times 10^{-2} \frac{m_h}{200 \text{ GeV}} \quad (18)$$

$$\sqrt{\left| A_{13}^{d\,2} + B_{13}^{d\,2} - \frac{14}{5} A_{13}^d B_{13}^d \right|} < 2 \times 10^{-2} \frac{m_h}{200 \text{ GeV}} \quad (19)$$

$$\sqrt{\left| A_{23}^{d\,2} + B_{23}^{d\,2} - \frac{14}{5} A_{23}^d B_{23}^d \right|} < 7 \times 10^{-2} \frac{m_h}{200 \text{ GeV}} \quad (20)$$

$$\sqrt{\left| A_{12}^{u\,2} + B_{12}^{u\,2} - \frac{14}{5} A_{12}^u B_{12}^u \right|} < 2 \times 10^{-2} \frac{m_h}{200 \text{ GeV}}. \quad (21)$$

These constraints are satisfied as long as $A_{12}^{u,d}$, $B_{12}^{u,d}$, $A_{13,23}^d$, $B_{13,23}^d$ are all $\mathcal{O}(\epsilon)$ or smaller. From the definition of the matrices $A^{u,d}$ and $B^{u,d}$ in eq. (13) it is easy to see that their off-diagonal elements ij are suppressed by powers of ϵ either if the corresponding mixing angle in the rotation matrices is suppressed ($V_{L,R\,ij} \sim \epsilon$) or if the corresponding a or b coefficients are universal ($a_i = a_j$ or $b_i = b_j$). As we will show in the following, for realistic assignments that reproduce the quark mass pattern, this is always the case, and the off-diagonal elements of the matrices $A^{u,d}$ and $B^{u,d}$ are indeed $\mathcal{O}(\epsilon)$ or smaller.

Constraints from $\Delta F = 1$ dipole processes like $b \rightarrow s\gamma$ are less restrictive, because the new contribution is suppressed by a loop factor and three Yukawa couplings¹. On the other hand, the CP-violating observable ϵ_K imposes a severe constraint on the new complex phases present in the effective theory,

$$\sqrt{\text{Im} \left(A_{12}^{d\,2} + B_{12}^{d\,2} - \frac{12}{5} A_{12}^d B_{12}^d \right)} < 4 \times 10^{-3} \frac{m_h}{200 \text{ GeV}}. \quad (22)$$

¹Similarly, Higgs-induced EDMs are small since they are also suppressed by a loop factor and three Yukawa couplings.

The Yukawa structures consistent with the size of the Cabibbo angle typically predict that either A_{12}^d or B_{12}^d is $\mathcal{O}(\epsilon)$ and therefore the constraint in eq. (22) is not naturally satisfied. However, in our approach, some of the c_{ij}^d coefficients in eq. (2) have to be somewhat smaller than one, and therefore it is possible to arrange their values such that the constraint in eq. (22) is satisfied. Since in the limit $c_{21}^d \ll 1$, the CP phases can be rotated away from the upper left 2×2 block of Y^d , choosing $|c_{21}^d| \leq 0.1$ allows us to satisfy the constraint from ϵ_K , without altering the quality of the fit to the quark masses and mixings. Alternatively, we can assume that the CP-violating phase entering eq. (22) (not related to the CKM phase) is about 0.1, rather than 1. This would be sufficient to satisfy the constraint. Finally, we mention that the CP-violating observable ϵ'_K imposes only a weaker constraint on $\text{Im}A_{12}^d$ and $\text{Im}B_{12}^d$.

Let us now give an example of a texture that leads to the mass and mixing pattern of eq. (5). Choosing

$$a = (1, 1, 0) , \quad b^d = (2, 1, 1) , \quad b^u = (2, 0, 0) , \quad (23)$$

we obtain²

$$Y^d \sim \begin{pmatrix} \epsilon^3 & \epsilon^2 & \epsilon^2 \\ \epsilon^3 & \epsilon^2 & \epsilon^2 \\ \epsilon^2 & \epsilon^1 & \epsilon^1 \end{pmatrix} , \quad Y^u \sim \begin{pmatrix} \epsilon^3 & \epsilon^1 & \epsilon^1 \\ \epsilon^3 & \epsilon^1 & \epsilon^1 \\ \epsilon^2 & \epsilon^0 & \epsilon^0 \end{pmatrix} . \quad (24)$$

The corresponding matrices $A^{u,d}$ and $B^{u,d}$, at leading order in ϵ , are given by

$$A^{u,d} = \begin{pmatrix} 1 & \epsilon^2 & \epsilon^1 \\ \epsilon^2 & 1 & \epsilon^1 \\ \epsilon^1 & \epsilon^1 & \epsilon^2 \end{pmatrix} , \quad B^d = \begin{pmatrix} 2 & \epsilon^1 & \epsilon^1 \\ \epsilon^1 & 1 & \epsilon^2 \\ \epsilon^1 & \epsilon^2 & 1 \end{pmatrix} , \quad B^u = \begin{pmatrix} 2 & \epsilon^2 & \epsilon^2 \\ \epsilon^2 & \epsilon^4 & \epsilon^4 \\ \epsilon^2 & \epsilon^4 & \epsilon^4 \end{pmatrix} . \quad (25)$$

Since the off-diagonal elements of $A^{u,d}$ and $B^{u,d}$ are $\mathcal{O}(\epsilon)$ or smaller, all $\Delta F = 2$ constraints are satisfied. The Higgs-mediated contribution to Δm_{B_d} is significant, nearly saturating the experimental value, as seen from eq. (19). On the other hand, since both A_{12}^u and B_{12}^u are $\mathcal{O}(\epsilon^2)$, the new contribution to Δm_D is rather small. We have checked that the texture defined by eq. (23) can reproduce the known values of quark masses and mixing, taking all coefficients $c_{ij}^{u,d}$ in the range 1/5 to 5.

²Equation (23) actually describes the unique factorizable and hierarchical solution leading to eq. (5). If we modify the assumption of eq. (5), taking $V_{us} \sim V_{cb} \sim \mathcal{O}(\epsilon^1)$ and $V_{ub} \sim \mathcal{O}(\epsilon^2)$, there is a unique solution with $a = (2, 1, 0)$, $b^d = (1, 1, 1)$, $b^u = (1, 0, 0)$ leading to

$$A^{u,d} = \begin{pmatrix} 2 & \epsilon^1 & \epsilon^2 \\ \epsilon^1 & 1 & \epsilon^1 \\ \epsilon^2 & \epsilon^1 & \epsilon^2 \end{pmatrix} , \quad B^u = \begin{pmatrix} 1 & \epsilon^1 & \epsilon^1 \\ \epsilon^1 & \epsilon^2 & \epsilon^2 \\ \epsilon^1 & \epsilon^2 & \epsilon^2 \end{pmatrix} ,$$

and B^d equal to the identity. Of course, more solutions exist if we drop the factorization hypothesis of eq. (7).

The extension of the Higgs-dependent Yukawa couplings hypothesis to the charged lepton sector is straightforward. The Higgs boson mediates interactions that violate lepton flavor and contribute to rare μ and τ decays and μ - e conversion processes. A simple estimate gives

$$BR(\mu \rightarrow e\gamma) \sim \frac{\alpha}{4\pi} \left(\frac{m_\mu^2 A_{12}^\ell}{m_h^2} \right)^2 = \left(\frac{200 \text{ GeV}}{m_h} \right)^4 \left(\frac{A_{12}^\ell}{1/60} \right)^2 1 \times 10^{-20} \quad (26)$$

$$BR(\mu \rightarrow eee) \sim \left(\frac{m_\mu m_e A_{12}^\ell}{m_h^2} \right)^2 = \left(\frac{200 \text{ GeV}}{m_h} \right)^4 \left(\frac{A_{12}^\ell}{1/60} \right)^2 5 \times 10^{-22} \quad (27)$$

$$BR(\tau \rightarrow \mu\gamma) \sim \frac{\alpha}{4\pi} \left(\frac{m_\tau^2 A_{23}^\ell}{m_h^2} \right)^2 = \left(\frac{200 \text{ GeV}}{m_h} \right)^4 \left(\frac{A_{23}^\ell}{1/60} \right)^2 1 \times 10^{-15} \quad (28)$$

$$BR(\tau \rightarrow \mu\mu\mu) \sim \left(\frac{m_\tau m_\mu A_{23}^\ell}{m_h^2} \right)^2 = \left(\frac{200 \text{ GeV}}{m_h} \right)^4 \left(\frac{A_{23}^\ell}{1/60} \right)^2 6 \times 10^{-15}. \quad (29)$$

The off-diagonal elements of the matrices A^ℓ and B^ℓ are not related to CKM angles. However, even if they have the same form as the corresponding elements in the down-quark sector and are of order ϵ , the predicted rates for lepton flavor violation satisfy the present experimental constraints and are too small to give a detectable signal in planned experiments.

We note that, with Higgs-dependent Yukawa couplings, CP violation is already possible in the system of two quarks, say the top and charm quarks. The analog of the Jarlskog invariant that controls CP violation in Higgs interactions is $\text{Im}[\text{Tr}(my^\dagger)^2]$. Unlike in the SM, it is not suppressed by a product of light quark masses. This opens up the possibility of EW baryogenesis which is to be studied elsewhere.

4 Higgs physics

One of the most striking features of our approach is that the Higgs branching ratios are drastically modified with respect to the SM predictions. The Higgs couplings to the weak gauge bosons and to the top quark (and consequently to gluons³ and photons) remain the same as those in the SM, but

$$\frac{\Gamma(h \rightarrow b\bar{b})}{\Gamma(h \rightarrow b\bar{b})_{SM}} = \frac{\Gamma(h \rightarrow c\bar{c})}{\Gamma(h \rightarrow c\bar{c})_{SM}} = \frac{\Gamma(h \rightarrow \tau^+\tau^-)}{\Gamma(h \rightarrow \tau^+\tau^-)_{SM}} = 9, \quad \frac{\Gamma(h \rightarrow \mu^+\mu^-)}{\Gamma(h \rightarrow \mu^+\mu^-)_{SM}} = 25. \quad (30)$$

The Higgs couplings to light quarks are enhanced even further, but they do not play any significant role in Higgs phenomenology. The enhancement of the coupling to μ can be very important in view of a possible muon collider, since the Higgs production rate is predicted to be a factor of 25 larger than that in the SM.

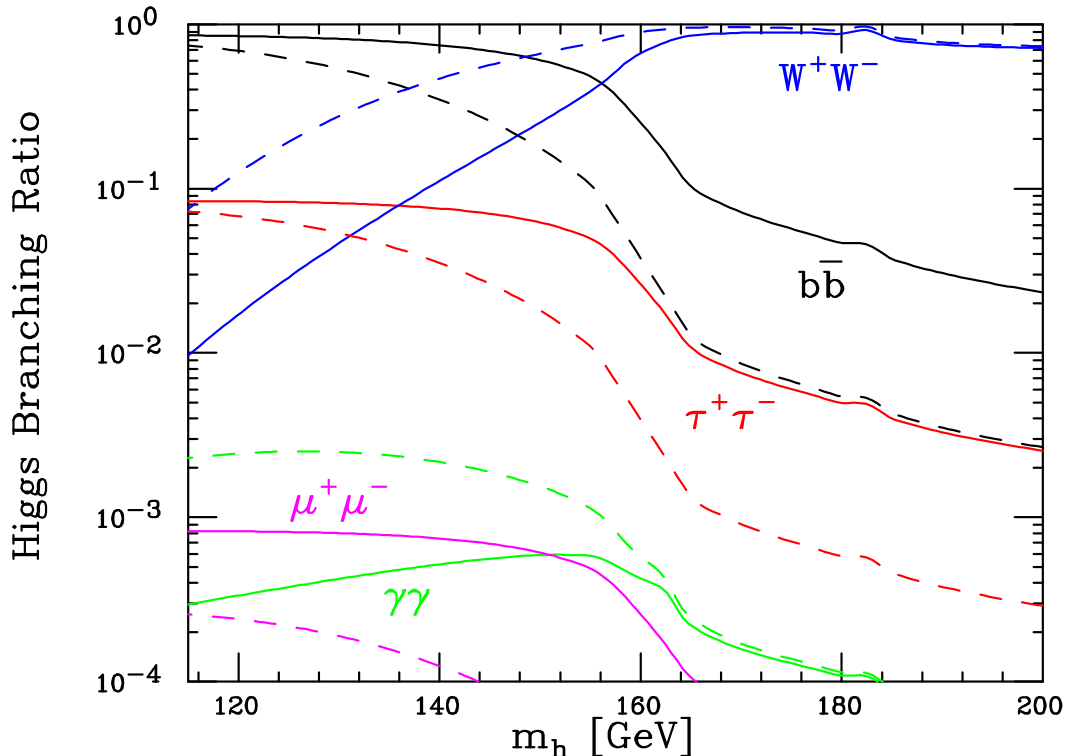


Figure 1: The various color lines show the Higgs branching ratios for different decay modes, with solid lines referring to the case of Higgs-dependent Yukawa couplings and dashed lines to the SM.

All main Higgs production processes at the LHC such as the gluon fusion, weak-boson fusion, Higgs-strahlung from the top or gauge boson are not affected, but the novelty lies in the Higgs decay. In fig. 1 we show the prediction for the Higgs branching ratios in the most important channels. At low m_h , there is an increase of $BR(h \rightarrow b\bar{b})$ and $BR(h \rightarrow \tau^+\tau^-)$ with respect to those in the SM, while, more importantly, there is also a significant reduction of $BR(h \rightarrow \gamma\gamma)$. Actually, the Higgs decay rate into muons becomes even larger than the one into photons, although its branching ratio remains smaller than 10^{-3} . At intermediate values of m_h , the main effect is an increase of the decay rate into fermions compared to that into WW and ZZ . As a result, $h \rightarrow WW$ becomes the leading decay mode only for $m_h > 156$ GeV, while in the SM this happens for $m_h > 136$ GeV.

Another peculiarity of our scenario concerns flavor-violating Higgs decay modes. The

³The modification of the Higgs-bottom coupling can affect the Higgs-gluon coupling, but the effect on the gluon-fusion rate is at most at the 10% level.

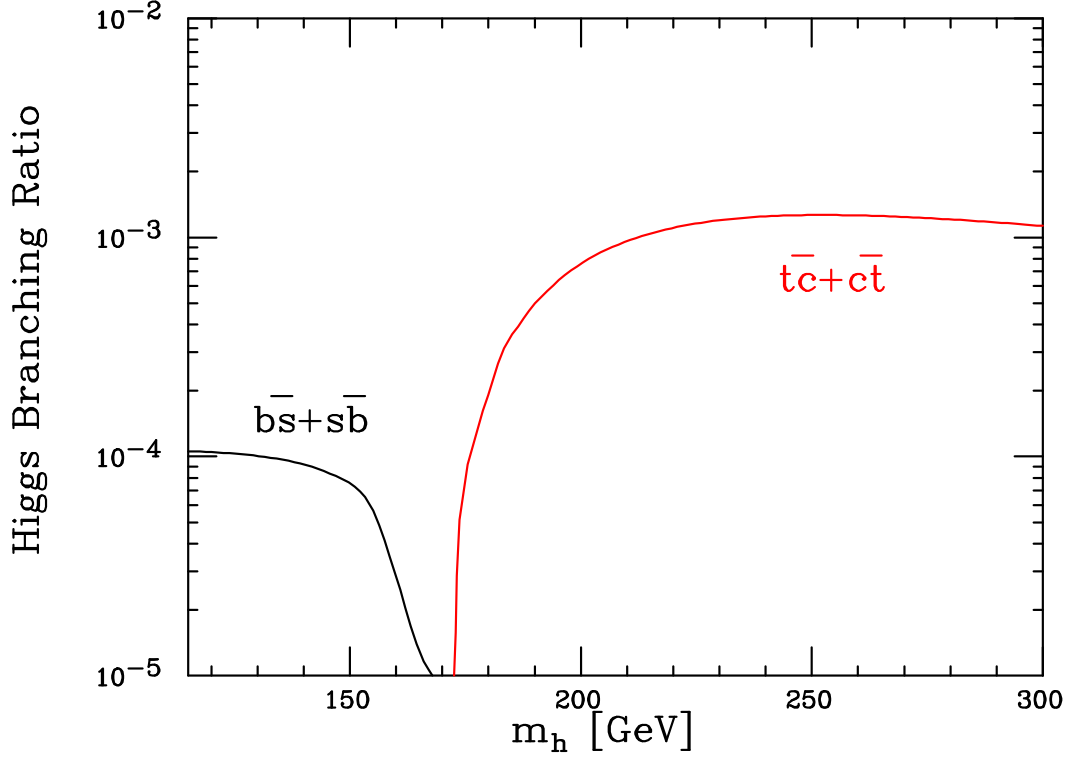


Figure 2: The Higgs branching ratios for flavor-violating decay modes, in the case of Higgs-dependent Yukawa couplings, taking $A_{23}^{u,d} = \epsilon = m_b/m_t$ and $B_{23}^{u,d} = 0$.

new Higgs interactions can be read off from eq. (11),

$$\mathcal{L}_h = -\frac{\sqrt{2}h}{v} \bar{u} [m^u (A^u P_L + B^u P_R) + (A^u P_R + B^u P_L) m^u] u + (u \leftrightarrow d), \quad (31)$$

where $P_{R,L} = (1 \pm \gamma_5)/2$ are the chiral projectors and $m^{u,d}$ are the diagonal quark mass matrices. From this we obtain

$$\frac{\Gamma(h \rightarrow b\bar{s} + \bar{b}s)}{\Gamma(h \rightarrow b\bar{b})_{SM}} = 4 \left(|A_{23}^d|^2 + |B_{23}^d|^2 \right), \quad (32)$$

$$\frac{\Gamma(h \rightarrow t\bar{c} + \bar{t}c)}{\Gamma(h \rightarrow b\bar{b})_{SM}} = 4 \left(|A_{23}^u|^2 + |B_{23}^u|^2 \right) \frac{m_t^2}{m_b^2} \left(1 - \frac{m_t^2}{m_h^2} \right)^2. \quad (33)$$

For $A_{23}^{u,d} = O(\epsilon)$, as given in eq. (25), the ratio in eq. (32) is about 10^{-3} . In eq. (33), the powers of ϵ from A_{23}^u are exactly compensated by the factor m_t^2/m_b^2 and the flavor-violating Higgs decay into the top quark, whenever kinematically allowed, has a rate of the order of the SM width for $h \rightarrow b\bar{b}$. We note however that the process $h \rightarrow tc$ is allowed only in the region of m_h where the dominant Higgs decay channel is $h \rightarrow WW$. In fig. 2 we show our prediction for the branching ratios of the flavor-violating Higgs modes, taking $A_{23}^{u,d} = \epsilon = m_b/m_t$ and

$B_{23}^{u,d} = 0$. In the range $200 \text{ GeV} \lesssim m_h \lesssim 300 \text{ GeV}$, $BR(h \rightarrow t\bar{c} + \bar{t}c)$ is typically of the order of 10^{-3} . For comparison, the corresponding branching ratio in the SM is of order 10^{-13} and that in two-Higgs doublet models is always less than 10^{-4} (see *e.g.* ref. [3]).

If the Higgs is lighter than the top quark, the interesting flavor-violating process is $t \rightarrow hc$. Its branching ratio is

$$BR(t \rightarrow hc) = \frac{2(|A_{23}^u|^2 + |B_{23}^u|^2) \left(1 - \frac{m_h^2}{m_t^2}\right)^2}{\left(1 - \frac{3m_W^4}{m_t^4} + \frac{2m_W^6}{m_t^6}\right)}. \quad (34)$$

For $A_{23}^{u,d} = O(\epsilon)$ and a light Higgs, $BR(t \rightarrow hc)$ is about 10^{-3} , which is well within the reach of the LHC, since experiments are expected to probe values down to 5×10^{-5} [4]. We note that the corresponding SM prediction is 6×10^{-15} , while type I and II two-Higgs doublet models give $BR(t \rightarrow hc) < 10^{-5}$ (see *e.g.* ref. [3]).

5 Example of TeV completion

In sect. 3 we have investigated flavor violation generated by Higgs exchange within the effective theory. However, since the mass M is around the TeV scale, higher-dimensional operators involving quarks and leptons, obtained from integrating out the heavy modes, can potentially be a dangerous source of FCNC. This issue cannot be addressed without specifying a particular model of flavor dynamics at the scale M . In this section we will present an example of a completion of the effective theory beyond the scale M , which does not lead to excessive flavor violation. Our example is not meant to describe a fully realistic theory of flavor, but only to illustrate how Higgs-dependent Yukawa couplings can be generated through couplings of quarks and leptons to some heavy fields.

Consider an extension of the SM with some heavy vectorlike Dirac fermions S^d, R^d having the gauge quantum numbers of the down-quark chiral components, i.e. $S^d \sim q_L$, $R^d \sim d_R$, with the interaction Lagrangian

$$\begin{aligned} -\mathcal{L} = & \left[\bar{q}_L \lambda_0^d d_R + \bar{q}_L \lambda_1^d R^d + \bar{S}^d \lambda_2^d d_R + \bar{S}^d \left(\lambda_3^{d\dagger} P_L + \lambda_4^d P_R \right) R^d \right] H + \text{h.c.} \\ & + \bar{S}^d m_S^d S^d + \bar{R}^d m_R^d R^d. \end{aligned} \quad (35)$$

Here λ_{0-4}^d and m_S, m_R are matrices in generation space and we have suppressed the flavor indices. We have also made the kinetic terms canonical and eliminated, by an appropriate basis transformation, possible mass terms which mix the light and heavy generations. We also introduce heavy fields S^u and R^u for the up-quarks in an analogous way. For simplicity,

we assume that the heavy fields from the up quark sector have negligible interactions with the down quark sector. The couplings λ_i^d and the mass scale $m_{S,R}/\lambda_i^d$ play the role of the parameters g_F and M that we have introduced in the effective theory in sect. 2.

Using their equations of motion, we integrate out the heavy states from eq. (35) and obtain modified kinetic terms for the quarks as well as the following Yukawa couplings

$$\begin{aligned} Y^d &= \lambda_0^d + \tilde{\lambda}_1^d \tilde{\lambda}_3^d \left(1 - \tilde{\lambda}_4^d \tilde{\lambda}_3^d\right)^{-1} \lambda_2^d \\ &= \lambda_0^d + \tilde{\lambda}_1^d \tilde{\lambda}_3^d \sum_{n=0}^{\infty} \left(\tilde{\lambda}_4^d \tilde{\lambda}_3^d\right)^n \lambda_2^d, \end{aligned} \quad (36)$$

where

$$\tilde{\lambda}_{1,4}^d = \lambda_{1,4}^d \frac{1}{m_R^d} H, \quad \tilde{\lambda}_3^d = \lambda_3^d \frac{1}{m_S^d} H^\dagger, \quad (37)$$

and analogously for the up-quark sector. Expanding the result in powers of $\tilde{\lambda}_4^d \tilde{\lambda}_3^d$, we obtain the desired Higgs-dependent Yukawa couplings as a series in ϵ .

Although generically the Yukawa couplings in eq. (36) do not have the factorizable form of eq. (7), with an appropriate choice of the flavor matrices λ_i , one can reproduce the texture of eq. (24). We find that in order to generate the texture in eq. (24), while avoiding excessive FCNC, one has to introduce 4 generations of fermions S and R . The simplest solution is to take equal masses for all S and R , flavor-conserving and universal couplings between the light and heavy fermions, while having all flavor violation reside in the interactions among the heavy fermions. For instance, we can choose

$$m_{S,R} = M \mathbb{1}_{4 \times 4}, \quad \lambda_1^{u,d} = \mathbb{1}_{3 \times 4}, \quad \lambda_2^{u,d} = \mathbb{1}_{4 \times 3}, \quad (38)$$

$$\begin{aligned} \lambda_3^d &= \begin{pmatrix} 0 & 0 & 0 & 1 \\ 0 & 0 & 0 & 1 \\ 0 & 1 & 1 & 0 \\ 1 & 0 & 0 & 0 \end{pmatrix}, \quad \lambda_4^d = \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 1 \\ 0 & 0 & 1 & 0 \end{pmatrix}, \\ \lambda_3^u &= \begin{pmatrix} 0 & 1 & 1 & 0 \\ 0 & 1 & 1 & 0 \\ 0 & 0 & 0 & 1 \\ 1 & 0 & 0 & 0 \end{pmatrix}, \quad \lambda_4^u = \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}, \end{aligned} \quad (39)$$

where $M \sim 1 - 2$ TeV, $(\mathbb{1}_{n \times m})_{ij}$ is an $n \times m$ matrix with elements 1 if $i = j$ and 0 otherwise, and ones and zeros in $\lambda_{3,4}^{u,d}$ are understood in the “texture” sense. The contribution to the quark wavefunctions from integrating out the heavy fields affects the Yukawa couplings at order ϵ^2 , which, for our purposes, can be neglected. Finally, the couplings of the light quarks

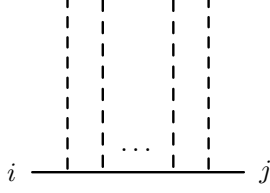


Figure 3: Fermion flavor violation via Higgs line insertions. The external legs of the fermion line represent quarks of generations i and j , while the internal propagators represent the heavy fermions R and S .

among themselves are

$$\lambda_0^d = 0, \quad \lambda_0^u = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 1 & 1 \end{pmatrix}. \quad (40)$$

Let us consider in more detail how FCNC are generated in this model. An important ingredient that allows us to suppress FCNC is the universality requirement in eq. (38), while eq. (39) represents a particular choice of λ_i which produces the desired Yukawa structure. Flavor violating operators are best studied in the weak basis rather than the mass eigenstate basis. Any Feynman diagram contributing to flavor violation is built out of fermion lines with insertions of Higgs (and possibly gauge boson) lines, see fig. 3. The Higgs lines either correspond to a Higgs VEV insertion or they are closed in loops (among themselves or against another fermion line). The resulting flavor structure can be studied order by order in the number of Higgs lines. Every two Higgs lines correspond to a factor of ϵ regardless of whether they represent the Higgs VEV insertions or they are closed in a loop (since ϵ is numerically close to the loop factor). Keeping the number of Higgs lines fixed, one sums over all possible λ_i matrices at the Higgs vertices.

The string of Higgs insertions in fig. 3 always starts and ends with the matrices λ_1 or λ_2 (or their conjugates), because we are interested in diagrams with external light quark lines. Since we require $\lambda_{1,2}$ to be unit matrices, non-trivial flavor structures can only arise if more than two Higgs lines are present. Consider the down sector. For 3 Higgs lines, the only possible structures are

$$\bar{d}_{L_i} \left(\lambda_1^d \lambda_3^d \lambda_2^d, \lambda_1^d \lambda_4^{d\dagger} \lambda_2^d \right)_{ij} d_{R_j}, \quad (41)$$

which could potentially generate processes like the $b \rightarrow s\gamma$ transition when a photon line is attached and two Higgs lines are closed in a loop. However, $\lambda_1^d \lambda_4^{d\dagger} \lambda_2^d$ is zero for our λ_i . Furthermore, $\lambda_1^d \lambda_3^d \lambda_2^d$ is proportional to the mass matrix for the down quarks (at order ϵ) in eq. (36), and thus the transformations that diagonalize the mass matrix will also diagonalize flavor along this fermion line. Thus no FCNC are generated at this level.

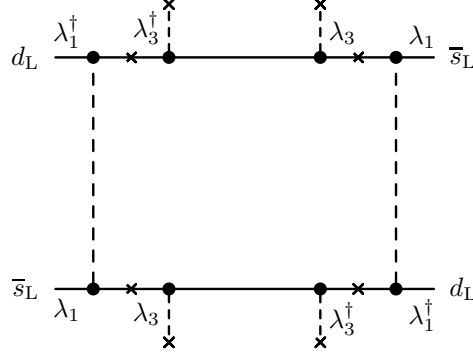


Figure 4: Example of a heavy quark contribution to the $K - \bar{K}$ mixing. Crosses indicate either a heavy quark mass or a Higgs VEV insertion.

The above cancellation is due to our universality requirement in eq. (38), which ensures that the matrices $\tilde{\lambda}_i$ defined in eq. (37) are proportional to the corresponding matrices λ_i . Effectively, there is an operative GIM mechanism which ensures that flavor violation is exactly rotated away, once we go to the mass eigenbasis. Loop corrections spoil the mass degeneracy of the heavy fermions, and we expect off-diagonal entries for the matrices $m_{S,R}$ of typical size $(\delta m_{S,R})_{ij} = (m_{S,R})_{ii} \lambda^2 / (16\pi^2)$. These effects are taken into account by a 5-Higgs insertion in the fermionic line, in which two Higgs vertices are closed in a loop by a Higgs propagator. The resulting contribution is sufficiently small and obeys the experimental limits on FCNC.

In the case of 4 Higgs lines, the transition along the fermion line is chirality conserving and a number of flavor objects appear. Let us consider the left-left transitions only, since the analysis for the right-right transitions is completely analogous. The possible structures along the fermionic string are

$$\bar{d}_{L_i} \left[\left(\lambda_1^d \lambda_3^d \lambda_4^d \lambda_1^{d\dagger} + \text{h.c.} \right), \lambda_1^d \lambda_3^d \lambda_3^{d\dagger} \lambda_1^{d\dagger}, \lambda_1^d \lambda_4^{d\dagger} \lambda_4^d \lambda_1^{d\dagger} \right]_{ij} d_{L_j}. \quad (42)$$

All of these combinations but the last one also contribute at tree level (order ϵ^2) to the quark kinetic terms $\bar{d}_{L_i} \not{\partial} d_{L_j}$. The term $\lambda_1^d \lambda_4^{d\dagger} \lambda_4^d \lambda_1^{d\dagger}$ is generated when all heavy quark propagators preserve chirality and thus leads to a higher-derivative operator. Although it does not affect the quark kinetic terms, it contributes, for instance, to the $\Delta F = 2$ box diagram obtained by coupling the Higgs propagators to another identical fermion line. This creates a mismatch between the flavor structure of the kinetic terms and that of the four-fermion operators. The transformations that diagonalize the kinetic terms do not then diagonalize flavor along the fermion line of the box diagram, leading to FCNC. In this case, the universality requirement in eq. (38) is not sufficient to guarantee a complete GIM mechanism. Thus FCNC appear at order ϵ^2 times a loop factor.

The most dangerous FCNC operator induced is $(\bar{s}_L \gamma_\mu d_L)^2$. An estimate of the diagram in fig. 4 yields

$$\mathcal{H}_{\text{eff}} \simeq \frac{\epsilon^2}{640\pi^2 M^2} (\bar{s}_L \gamma_\mu d_L)^2. \quad (43)$$

This induces Δm_K of order 10^{-16} GeV which is below the experimental limit. Other meson mixing constraints are weaker. On the other hand ϵ_K is quite restrictive in general and arbitrary CP phases typically overproduce it by an order of magnitude. This constraint, however, depends on how CP violating phases enter the flavor structures. For example, if CP violation is mostly due to the up quark sector interactions, the ϵ_K constraint is not significant. We also note that the heavy quark contribution to the real and imaginary parts of the $K - \bar{K}$ mixing can be suppressed further for other choices of $\lambda_{3,4}^d$. In particular, we have found 5×5 textures that induce $K - \bar{K}$ mixing only at order ϵ^3 .

Interactions with the heavy quarks also induce non-trivial kinetic structures of the form of the second term in eq. (9). To lowest order these terms are

$$\frac{(\lambda_1^d \lambda_1^{d\dagger})_{ij}}{M^2} \overline{(H^\dagger q_{Li})} i \not{D} (H^\dagger q_{Lj}) + \frac{(\lambda_1^u \lambda_1^{u\dagger})_{ij}}{M^2} \overline{(H q_{Li})} i \not{D} (H q_{Lj}), \quad (44)$$

and analogous terms are generated for u_R and d_R . Here $H q_L \equiv \epsilon^{ab} H_a q_{Lb}$ with a, b being the SU(2) indices. Since $\not{D}(H^\dagger q_{Lj})$ and $\not{D}(H q_{Lj})$ have a different gauge structure compared to that of $\not{D} q_{Lj}$, we find a modification of the Z -fermion couplings. After canonically redefining the fields, the tree-level couplings $g_{L,R}$ of the left and right-handed quarks to the Z are given by

$$\begin{aligned} g_L &= I_3 - \sin^2 \theta_W Q + \delta g_L & \delta g_L &= -I_3 \frac{\lambda_1^{u,d} \lambda_1^{u,d\dagger} v^2}{M^2}, \\ g_R &= -\sin^2 \theta_W Q + \delta g_R & \delta g_R &= I_3 \frac{\lambda_2^{u,d\dagger} \lambda_2^{u,d} v^2}{M^2}. \end{aligned} \quad (45)$$

Here I_3 is the third isospin component ($I_3 = \pm 1/2$ for up and down quarks, respectively) and Q is the fermion electric charge. For universal $\lambda_{1,2}^{u,d}$ and $m_{R,S}$, as in eq. (38), the correction affects only $g_A \equiv g_L - g_R$ and not $g_V \equiv g_L + g_R$, and we find $\delta g_A = O(\epsilon)$. For $\epsilon \sim 10^{-2}$, which is appropriate to reproduce the hierarchical pattern of eq. (5), δg_A is not inconsistent with experimental data on g_A for charm and bottom quarks, since the 2σ errors are $\Delta g_A^c = 1.06 \times 10^{-2}$ and $\Delta g_A^b = 1.02 \times 10^{-2}$ [5]. Actually, if we simultaneously include the corrections to all quark couplings, we find that the strongest constraint is imposed by $R_b = 0.21629 \pm 0.00066$ [5] which allows for a 0.6% deviation at the 2σ level. The corrections from eq. (45) induce $\delta R_b/R_b = 0.28\epsilon$, which is 0.5% for $\epsilon = 1/60$. Note that there is freedom in the model to reduce the contributions to $\delta g_{L,R}$ while keeping the quark masses fixed, by decreasing the values of $\lambda_{1,2,4}$ and increasing λ_3 . This rescaling is limited by the requirement of maintaining

perturbative couplings. Finally, similar corrections to the W -fermion couplings are relatively weakly constrained and do not lead to further bounds.

The situation is more problematic when we try to extend the model to leptons, since g_A^ℓ has been measured with the precision of 10^{-3} at $2\text{-}\sigma$ [5] which requires an order of magnitude suppression of our δg_A^ℓ . Once ϵ is fixed by m_τ to be 10^{-2} , the correction δg_A^ℓ can be reduced by decreasing $\lambda_{1,2}$. To keep m_τ intact, one has to increase $\lambda_3 \simeq m_\tau/(v\delta g_A)$ to strong-coupling values of around 10, which renders our calculations unreliable. Another possibility is to produce m_τ with the direct coupling λ_0 and generate only m_μ and m_e through the mixing with heavy fermions, but this would go against the spirit of our approach. Finally, δg_A^ℓ can be reduced by a flavor-universal contribution of additional particles, which do not affect the flavor structures. We thus conclude that the toy model presented in this section cannot be directly extended to leptons and some new interactions are necessary to generate m_τ without inducing too large $\delta g_{L,R}^\ell$.

Let us return to the quark sector. After we integrate out at tree level the heavy fermions in the down-left sector, the induced interactions of order ϵ^2 are

$$i\bar{q}_L\tilde{\lambda}_1^d\left[\left(\tilde{\lambda}_3^d\tilde{\lambda}_4^d\mathcal{D}+\text{h.c.}\right)+\tilde{\lambda}_3^d\mathcal{D}\tilde{\lambda}_3^{d\dagger}-\mathcal{D}\tilde{\lambda}_4^{d\dagger}\frac{\mathcal{D}}{m_S^2}\tilde{\lambda}_4^d\mathcal{D}\right]\tilde{\lambda}_1^{d\dagger}q_L, \quad (46)$$

which correspond to the flavor structures shown in eq. (42). The second interaction contains terms of the form $\mathcal{D}(H^\dagger H q_{Lj})$, which have the same gauge structure as $\mathcal{D}q_{Lj}$ and therefore lead to an overall rescaling of the kinetic terms $\bar{q}_{Li}\mathcal{D}q_{Lj}$. On the other hand, the first interaction in eq. (46) is of the form

$$\frac{(\lambda_1^d\lambda_3^d\lambda_4^d\lambda_1^{d\dagger})_{ij}}{M^4}\bar{q}_{Li}HH^\dagger Hi\mathcal{D}(H^\dagger q_{Lj})+\text{h.c.}, \quad (47)$$

involving the different gauge structure $\mathcal{D}(H^\dagger q_{Lj})$ and thus leading to modified Z -fermion couplings. In particular, we find flavor violating $Z\bar{d}_{Li}d_{Lj}$ and $Z\bar{u}_{Li}u_{Lj}$ vertices

$$(\delta g_L^d)_{ij}=\frac{\epsilon^2}{2}(\lambda_1^d\lambda_3^d\lambda_4^d\lambda_1^{d\dagger}+\text{h.c.})_{ij}, \quad (\delta g_L^u)_{ij}=-\frac{\epsilon^2}{2}(\lambda_1^u\lambda_3^u\lambda_4^u\lambda_1^{u\dagger}+\text{h.c.})_{ij}. \quad (48)$$

In the right-handed sector, the corresponding flavor violating structures are $\lambda_2^{d\dagger}\lambda_4^d\lambda_3^d\lambda_2^d$ and $\lambda_2^{u\dagger}\lambda_4^u\lambda_3^u\lambda_2^u$. Such couplings are strongly constrained by flavor physics. The $Z\bar{s}d$ vertex is constrained most severely by $K^+\rightarrow\pi^+\nu\bar{\nu}$ and $K_L\rightarrow\mu^+\mu^-$ processes such that the effective coupling $Z\bar{s}_Ld_L$ must be $\lesssim 7\times 10^{-6}$ [6]. The $Z\bar{b}s$, $Z\bar{b}d$ vertices are constrained at the level $\text{few}\times 10^{-4}$ by $B_{s,d}\rightarrow\mu^+\mu^-$ and $b\rightarrow sl^+l^-$ processes [7]. This implies that the $Z\bar{s}d$ vertex should appear at the ϵ^3 level, while the $Z\bar{b}s$, $Z\bar{b}d$ vertices are allowed at the ϵ^2 order. This is indeed what we have in our model as $\lambda_1\lambda_3\lambda_4\lambda_1^\dagger$ and $\lambda_2^\dagger\lambda_4\lambda_3\lambda_2$ in both up- and down-sectors

have a zero 2×2 block in the upper left corner such that the (12) transition only appears at order ϵ^3 , but the (13) and (23) transitions already exist at order ϵ^2 . Therefore the flavor physics bounds are satisfied.

We remark that loops with heavy quarks which produce dipole interactions $\frac{1}{v} \bar{d}_{Li} \sigma_{\mu\nu} d_{Rj} F^{\mu\nu}$ lead to flavor violation in the down sector at order ϵ^3 (or, more precisely, ϵ^2 times a loop factor) in our model. This is because $\lambda_1^d \lambda_4^{d\dagger} \lambda_2^d = 0$ and the lowest order operators vanish. The resulting contribution to $\text{BR}(b \rightarrow s\gamma)$ is small. Finally, the induced quark EDMs depend strongly on how CP violation is implemented in the model. In particular, if the CKM phase is due to non-removable CP phases in λ_0 of eq. (40) while the other flavor objects are real in that basis, the EDM constraints are insignificant.

To conclude, the above model provides an example of how TeV scale new physics can generate Higgs-dependent structures in the Yukawa couplings. Although this model as it stands is viable only for quarks, while for leptons additional flavor-universal interactions are required, it shows that it is possible to induce these flavor structures without entailing excessive FCNC.

6 Conclusions

We have studied the possibility that the SM Yukawa couplings are functions of the Higgs field. This can explain the hierarchical structure of the fermion masses and mixing angles in terms of powers of the ratio between the Higgs VEV and a new mass scale. We find that this mass scale, characterizing the dynamics of flavor physics, is determined to be roughly $\sqrt{m_t/m_b} v \simeq 1\text{--}2$ TeV.

An immediate consequence of this approach is that the Higgs boson couplings to fermions are drastically modified. The Higgs decay widths into bottom quarks and into τ 's are 9 times larger than those in the SM, while the one into muons is larger by a factor of 25. The prediction for the Higgs branching ratios is shown in fig. 1. Furthermore, the Higgs couplings violate flavor and this results in observable rates for either $h \rightarrow t\bar{c}$ or $t \rightarrow hc$, whose branching fractions are expected to be of order 10^{-3} , when kinematically accessible. Tree-level Higgs exchange also contributes to various FCNC processes, but the effects are consistent with the present constraints. The most significant contribution is in the $B\text{--}\bar{B}$ system and can be very close to the current experimental sensitivity. The new flavor dynamics at the TeV scale is within the reach of LHC experiments.

We have also presented an example of a possible TeV scale completion of our effective

theory. This model involves heavy vector-like quarks with flavor-universal masses which interact with ordinary quarks in a flavor conserving way (at leading order). Due to this universality, dangerous FCNC operators are sufficiently suppressed, while the correct fermion masses and mixings can be generated.

Acknowledgements. We are grateful to G. Degrassi, S. Jäger, H.D. Kim, R. Rattazzi, M. Ratz, P. Slavich and T. Takeuchi for discussions. We would also like to thank K. Babu for pointing out ref. [1].

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